

Spontaneous Casimir Effect

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Abstract

The dynamical Casimir effect predicts that vacuum amplification effects, resulting in the creation of real particles out of vacuum fluctuations, are induced by rapidly modulating the boundary conditions of a quantum field. Here we show that a spontaneous release of virtual photons from the quantum vacuum can occur in a quantum optical system in the ultrastrong coupling regime. In contrast to the dynamical Casimir effect and other pair creation mechanisms, this phenomenon does not require external forces or time dependent parameters. It resembles the production of particles during the early universe expansion induced by the decay of a false vacuum according to inflationary cosmology

The dynamical Casimir effect [3] has been recently experimentally realized in a superconducting circuit by modulating the inductance of a quantum interference device at high frequencies [4]. Other proposed vacuum amplification mechanisms [1], as the Schwinger process [12] and the Hawking radiation [13], require the presence of huge external fields or, as the Unruh effect, the presence of a rapidly nonuniformly accelerating observer [14], and have yet to be observed. The here described spontaneous release of virtual photon pairs in the absence of any external drive or observer acceleration resembles the particle production during the universe expansion predicted by modern inflationary cosmology. According to this theory, a scalar field starting in the false vacuum state eventually decays, and the energy that had been locked in it is released to form a hot, uniform soup of particles, which is the assumed starting point of the traditional big bang theory [11].

The Hamiltonian of a realistic atom-cavity system contains so-called counter-rotating terms allowing the simultaneous creation or annihilation of an excitation in both atom and cavity mode. These terms can be safely neglected for small coupling rates Ω_R (rotating-wave approximation). However, when Ω_R becomes comparable to the cavity resonance frequency of the emitter or the resonance frequency of the cavity mode, the counter-rotating terms are expected to manifest, giving rise to exciting effects in cavity QED [6, 15, 16]. This ultrastrong-coupling regime is difficult to reach in quantum-optical cavity QED, but was recently realized in a variety of solid-state quantum systems [5–10]. Such regime is challenging from a theoretical point of view as the total number of excitations in the cavity-emitter system is not preserved, even though its parity is [16].

It has been shown that, in the ultrastrong coupling regime, the quantum optical master equation fails to provide the correct description of the system's interaction with reservoirs [17]. Moreover quantum optical normal order correlation functions fail to describe photodetection experiments for such systems [18, 19]. Specifically, for a single mode resonator, the photon rate that can be detected by a photoabsorber is no more proportional to $\langle a^\dagger(t)a(t) \rangle$ (where a and a^\dagger are the photon destruction and creation operators) but to $\langle X^-(t)X^+(t) \rangle$, where $X^+(t)$ is the positive frequency component of the quadrature operator $X(t) = a(t) + a^\dagger(t)$ [19]. The most puzzling property of these systems is that their ground state is a squeezed vacuum containing correlated pairs of cavity photons [20]. The photon pairs in the ground state $|\tilde{0}\rangle$ are, however, virtual and cannot be detected [21], being $\langle \tilde{0}|X^-(t)X^+(t)|\tilde{0}\rangle = 0$ [19]. Otherwise, an observation of a stream of photons from such

a system in its ground state would give rise to *perpetuum mobile* behaviors. Nevertheless, we show that such virtual photon pairs, dragged by spontaneous decay processes reaching or starting from states of the quantum emitter not coupled with the resonator, may be spontaneously released and directly observed. A detailed analysis, carried out below, shows that such puzzling phenomenon is traceable to a general feature of open quantum systems: spontaneous transitions induced by a reservoir occur among eigenstates of the total Hamiltonian (environment induced superselection of energy eigenstates [22]). Such general feature, in presence of a ground state with virtual excitations (or also excited states that are not orthogonal to the zero-excitation state), induces the spontaneous Casimir effect here presented. A mechanism for the generation of quantum vacuum radiation, based on the presence of counter-rotating terms in the light-matter interaction Hamiltonian, has already been proposed [15, 18, 23]. Nevertheless these vacuum amplification proposals, as the dynamical Casimir effect, require a fast time-modulation or the sudden switch on or off of the vacuum Rabi frequency.

The most promising candidates for an experimental realization of the proposed setting are superconducting quantum circuits [24] and intersubband quantum well polaritons [5]. In particular, ultrastrong coupling and cascade multi-level structure (the two required features) have been experimentally demonstrated in separate experiments with quantum circuits [6, 24, 26] and in a single device with multiple quantum wells [5].

In the absence of losses, the quantum system depicted in Fig. 1a,b is described by the total Hamiltonian $H = H_s + H_{\text{Rabi}}$ with

$$H_{\text{Rabi}} = \omega_0 a^\dagger a + \sum_{\alpha=\text{g,e}} \omega_\alpha \sigma_{\alpha\alpha} + \Omega_{\text{R}}(a + a^\dagger)(\sigma_{\text{eg}} + \sigma_{\text{ge}}), \quad (1)$$

and $H_s = \omega_s \sigma_{ss}$, where $\omega_{0(\alpha)}$ ($\alpha = \text{s, g, e}$) are the bare energies of the cavity mode and of the atomic-like levels, and $\sigma_{\alpha\beta}$ describes the transition operators involving the levels of the quantum emitter. It is useful to label the eigenstates $|j\rangle$ with increasing eigenenergies Ω_j (j integer). Actually the Hamiltonian is block-diagonal and its eigenstates can be separated into (i) a non-interacting sector $|s, n\rangle$ with energy $\omega_s + n\omega_0$, where n labels the cavity photon number; and (ii) the dressed atom-cavity states resulting from the diagonalization of H_{Rabi} : $|\tilde{j}\rangle$ with energy $\omega_{\tilde{j}}$. The ground state of this sector can be expanded as

$$|\tilde{0}\rangle = \sum_{k=0}^{\infty} (c_{\text{g},2k}^{\tilde{0}} |\text{g}, 2k\rangle + c_{\text{e},2k+1}^{\tilde{0}} |\text{e}, 2k+1\rangle). \quad (2)$$

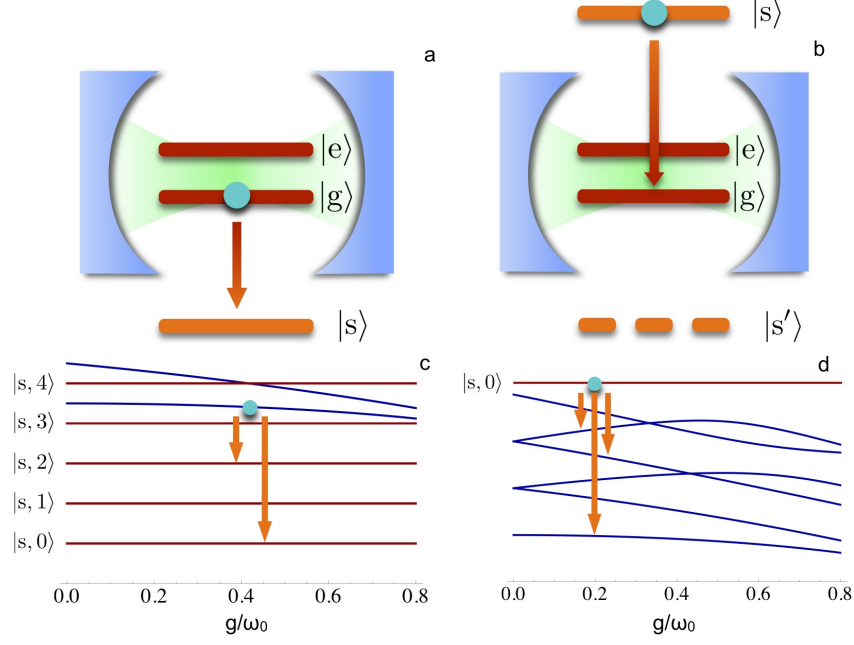


FIG. 1: (color online) **Sketch of the system under consideration.** Two-levels of a single quantum emitter are ultrastrongly coupled to a single cavity mode. In the upper panels spontaneous emission processes for two different configurations are sketched neglecting the ultrastrong coupling. Panels c and d show the involved total energy levels as a function of the coupling Ω_R/ω_0 and a sketch of the allowed spontaneous transitions in the ultrastrong coupling regime, corresponding to the configurations in panels a and b respectively. The bending lines as a function of Ω_R/ω_0 describes the dressed energy levels $\omega_{\tilde{j}}$.

For $\Omega_R/\omega_0 \ll 1$, the Jaynes-Cummings (JC) model is recovered and only $c_{g,0}^{\tilde{0}} \neq 0$. For increasing couplings, e.g. $\Omega_R/\omega_0 \approx 0.1$, the one photon-pair amplitude $c_{g,2}^{\tilde{0}} \sim \Omega_R/\omega_0$ becomes non-negligible. Analogously the excited states can be expanded as

$$|\tilde{j}\rangle = \sum_{k=0}^{\infty} (c_{g,k}^{\tilde{j}} |g, k\rangle + c_{e,k}^{\tilde{j}} |e, k\rangle). \quad (3)$$

In contrast to JC model, they (only those with the same parity of the state $|g, 0\rangle$) are not orthogonal to $|g, 0\rangle$ (being $c_{g,0}^{\tilde{j}} \neq 0$), hence they have a nonzero probability to have zero bare excitations. All subsequent calculations are performed at zero detuning: $\omega_0 - \omega_{eg} = 0$, where $\omega_{ij} = \omega_i - \omega_j$.

In the first of the two configurations (see Fig. 1a,c) that we consider, the system is initially prepared, e.g. by adiabatic excitation, in the lowest energy dressed state $|\tilde{0}\rangle$ of the

system. Although in the ultrastrong coupling regime this state contains virtual cavity photons ($\langle \tilde{0} | a^\dagger a | \tilde{0} \rangle \neq 0$), it is a vacuum state for the physical cavity photons that can be detected: $X^+ | \tilde{0} \rangle = 0$ and hence $\langle \tilde{0} | X^- X^+ | \tilde{0} \rangle = 0$. Since $| \tilde{0} \rangle$ is not the lowest energy vacuum, it can be considered as the photonic false vacuum, since in this case the state $| s, n = 0 \rangle$ is the lowest energy (true) vacuum state. Interesting theoretical studies of quantum dynamics in these cavity QED systems with [25] and without [27] the rotating-wave approximation recently appeared. Panel 1c displays the lowest energy levels (the eigenvalues of H : $\omega_s + n\omega_0$ and $\omega_{\tilde{j}}$) as a function of Ω_R . Let us discuss the spontaneous decay of the initial state $| I_a \rangle = | \tilde{0} \rangle$, keeping in mind that spontaneous transitions induced by a reservoir occur among eigenstates of the total Hamiltonian (environment induced superselection of energy eigenstates [22]). For zero or small coupling rates Ω_R , the initial state $| \tilde{0} \rangle$ reduces to $| g, 0 \rangle$ and standard spontaneous emission of a photon (at energy ω_{gs}) in the external electromagnetic modes, associated with the emitter transition $| g \rangle \rightarrow \sigma_{sg} | g \rangle = | s \rangle$, occurs at a rate γ_{gs} (fixed by the dipole moment of the transition). In the ultrastrong coupling regime, the initial state is $| I_a \rangle = | \tilde{0} \rangle$ and possible final states are $| F_i \rangle = | s, 2i \rangle$, where $\langle F_i | \sigma_{sg} | I_a \rangle = c_{g,2i}^{\tilde{0}}$. For $i = 0$, the final state contains no cavity photons as in ordinary spontaneous emission. For coupling rates $\Omega_R/\omega_0 < 1$, the contribution with $i = 1$ provides the next dominant term. Hence the spontaneous emission of a photon not in the cavity mode at a rate γ_{gs} comes together with a flux of cavity photon pairs at a rate $\approx \gamma_{gs} | c_{g,2}^{\tilde{0}} |^2$.

In the second configuration, shown in Fig. 1b, the system is initially prepared into the cavity-uncoupled level $| I_b \rangle = | s, 0 \rangle$, which is at higher energy with respect to state $| e \rangle$. In this case, the initial state can be prepared exciting the matter system from a ground state $| s' \rangle$, without involving the emitter states coupled with the resonator. However such setting works even in the absence of $| s' \rangle$. We consider the case where spontaneous emission induces the following transition: $| s \rangle \rightarrow | g \rangle$, as the transition $| s \rangle \rightarrow | e \rangle$ is forbidden. For small coupling rates $\Omega_R/\omega_0 \ll 1$, standard spontaneous emission of a photon at energy ω_{sg} directly in the external electromagnetic modes at a rate γ_{sg} would be observed. In the ultrastrong coupling regime possible final states are $| F_{\tilde{j}} \rangle = | \tilde{j} \rangle$, where $\langle F_{\tilde{j}} | \sigma_{gs} | I_b \rangle = c_{g,0}^{\tilde{j}}$. The coefficient $c_{g,0}^{\tilde{j}}$ describes the probability amplitude that the state $| \tilde{j} \rangle$ has zero cavity photons. For negligible couplings, only $c_{g,0}^{\tilde{0}} \neq 0$, thus the only possible final state would be $| \tilde{0} \rangle$. In this case no cavity photons would be observed, since $\langle \tilde{0} | X^- X^+ | \tilde{0} \rangle = 0$. In the ultrastrong coupling regime, $c_{g,0}^{\tilde{j}} \neq 0$ for even excited states (with the same parity as $| g, 0 \rangle$),

hence spontaneous transitions $|s, 0\rangle \rightarrow |F_{\tilde{j}}\rangle$ can occur even for $\tilde{j} \neq \tilde{0}$. In particular, the next largest coefficients are those for $\tilde{j} = 3, 4$. These excited dressed states emit physical photon pairs ($\langle \tilde{3}|X^-X^+|\tilde{3}\rangle \neq 0$). It is amazing and paradoxical that cavity photon-pairs are emitted only because there are excited states of the cavity-matter system which have a nonzero probability of having zero cavity photons!

For describing a realistic system, the cavity and atom dissipation channels need to be taken into account. Yet, owing to the very high ratio Ω_R/ω_0 , the description offered by the standard quantum optical master equation can break down, for example producing spurious qubit flipping or photon generation, even at zero temperature [17]. Following Refs. [17], we write the Hamiltonian in a basis formed by the eigenstates of H to describe the dissipative processes. We choose a $T = 0$ temperature environment. Yet generalization to $T \neq 0$ environments is straightforward. We thus arrive at the master equation [17, 25], $\dot{\rho}(t) = i[\rho(t), H] + \sum_c \mathcal{L}_c \rho(t)$, where \mathcal{L}_c is a Liouvillian superoperator describing the cavity ($c = 0$) and the material system losses $c = e \rightarrow g$, and $g \rightarrow s$ (configuration a) or $s \rightarrow g$ (b) (see Methods for details).

According to the input-output relations, the emitted photon flux out of the resonator is proportional to the mean value of intracavity physical photons: $\gamma_0 \langle X^-X^+ \rangle$. The results of a full numerical demonstration including the cavity and the emitter losses is shown in Fig. 2. Figure 2a displays the numerically calculated time evolution of $\langle X^-X^+ \rangle = \text{Tr}[X^-X^+\rho(t)]$ for configuration (a) (see Fig. 1a and c) and for different coupling strengths $\Omega_R/\omega_0 = 0.3$ (red line), 0.4 (blue), 0.6 (black). The signal rapidly grows and reaches a maximum value before decaying exponentially due to cavity losses. The signal increases with increasing Ω_R/ω_0 , as a consequence of the buildup of $c_{g,2}^{\tilde{0}}$. Calculations have been performed at zero detuning and by using $\gamma_{eg} = \gamma_0 = \gamma_{gs} = 2 \times 10^{-2}\Omega_R$. Panel 2b shows calculations of $\langle X^-X^+ \rangle$ for different spontaneous emission decay rates ($\gamma_{gs}/\Omega_R = 10^{-2}, 1.5 \times 10^{-2}, 3 \times 10^{-2}, 4 \times 10^{-2}$) obtained by artificially dropping cavity losses ($\gamma_0 = 0$). We also used a coupling rate $\Omega_R/\omega_0 = 0.6$ and $\gamma_{eg} = \gamma_{gs} = 2 \times 10^{-2}\Omega_R$. In the absence of cavity losses, the mean photon number reaches a maximum value which does not depend on γ_{gs} . This result puts forward that the phenomenon here investigated is intrinsically different from the dynamical Casimir effect where the emitted photon rate depends strongly on the modulation frequency. The inset in Fig. 2b displays the maximum cavity mean photon number $\langle X^-X^+ \rangle_{\text{Max}}$ as a function of the energy difference $\omega_{\tilde{0}}$ obtained using $\Omega_R/\omega_0 = 0.8$ (other parameters are the same as

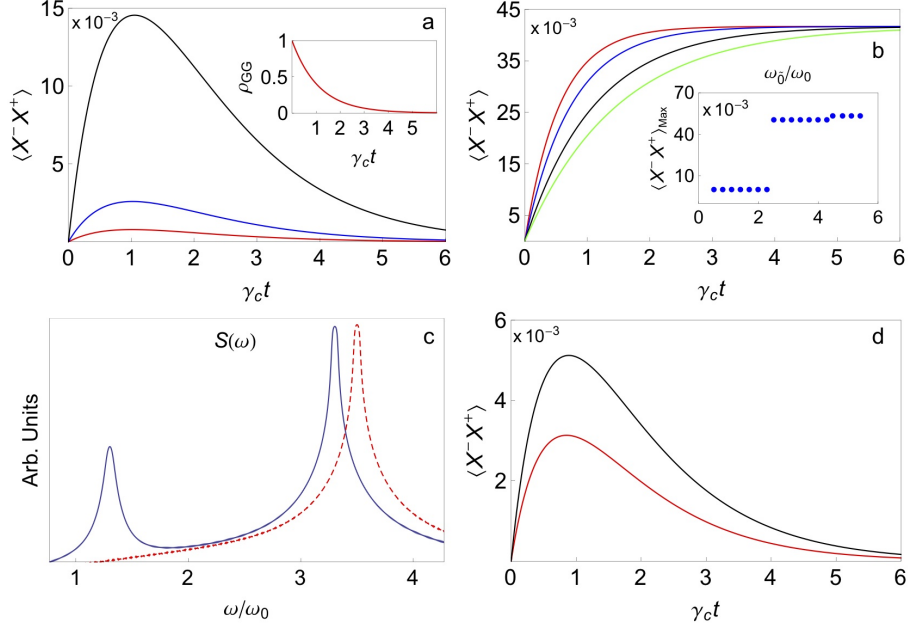


FIG. 2: (color online) **Dynamics of released photons and spontaneous emission spectra.**

a: Time evolution of $\langle X^- X^+ \rangle$ for configuration (a) (see Fig. 1a) and for different coupling strengths (see text). b: $\langle X^- X^+ \rangle$ for different spontaneous emission decay rates γ_{gs} obtained by artificially dropping cavity losses. The inset in Fig. 2b displays the maximum cavity mean photon number $\langle X^- X^+ \rangle_{\text{Max}}$ as a function of the energy difference $\omega_{\tilde{0}}$. c: Spectrum $S(\omega)$ of spontaneously emitted photons. d: Time evolution of $\langle X^- X^+ \rangle$ for the configuration (b) sketched in Fig. 1b.

those used for Fig. 2a). Photon pairs are released only when $\omega_{\tilde{0}} > \omega_s + 2\omega_0$ because the spontaneous transition $|\tilde{0}\rangle \rightarrow |s, 2\rangle$ occurs only when the energy of the final state is below that of the initial state (see Fig. 1c). Increasing $\omega_{\tilde{0}}$ a second rung corresponding to a small increase of $\langle X^- X^+ \rangle_{\text{Max}}$ can be observed. It comes from the activation of the transition $|\tilde{0}\rangle \rightarrow |s, 4\rangle$.

The release of virtual photon pairs, present in the photonic false vacuum $|\tilde{0}\rangle$, satisfies energy conservation. The energy of the radiated photons, induced by the spontaneous decay of the false vacuum, comes at the expense of the energy of the quanta emitted into the reservoir (e.g. the spontaneously emitted photons in the external electromagnetic modes). In the dynamical Casimir effect, the energy of radiated photons comes at the expense of the mechanical energy of the moving mirror, which then experiences a friction force by the quantum vacuum via the so-called back-reaction effect [27, 28]. In the present case, back-

reaction effects can be investigated by calculating the emission spectrum of spontaneously emitted photons: $S(\omega) = 1/(2\pi) \int_{-\infty}^{\infty} dt \int_{-\infty}^{\infty} dt' \langle \sigma^-(t) \sigma^+(t') \rangle e^{-i\omega(t-t')}$, where σ^\pm are the frequency positive (negative) components of the polarization operator $\sigma_{gs} + \sigma_{sg}$. The spectrum $S(\omega)$ of spontaneously emitted photons, calculated for $\Omega_R/\omega_0 = 0.6$ (other parameters are the same as those used for Fig. 2a), is displayed in Fig. 2c. In the absence of interaction (dashed line), the spectrum consists of a single Lorentzian peak centered at energy ω_{gs} . In the ultrastrong coupling regime the main peak is red-shifted at energy $\omega_{\tilde{0}s}$ and a second peak at lower energy, centered at $\omega_{\tilde{0}s} - 2\omega_0$ (see Fig. 1c) appears. It corresponds to spontaneously emitted photons of lower energy associated to the transition $|\tilde{0}\rangle \rightarrow |s, 2\rangle$.

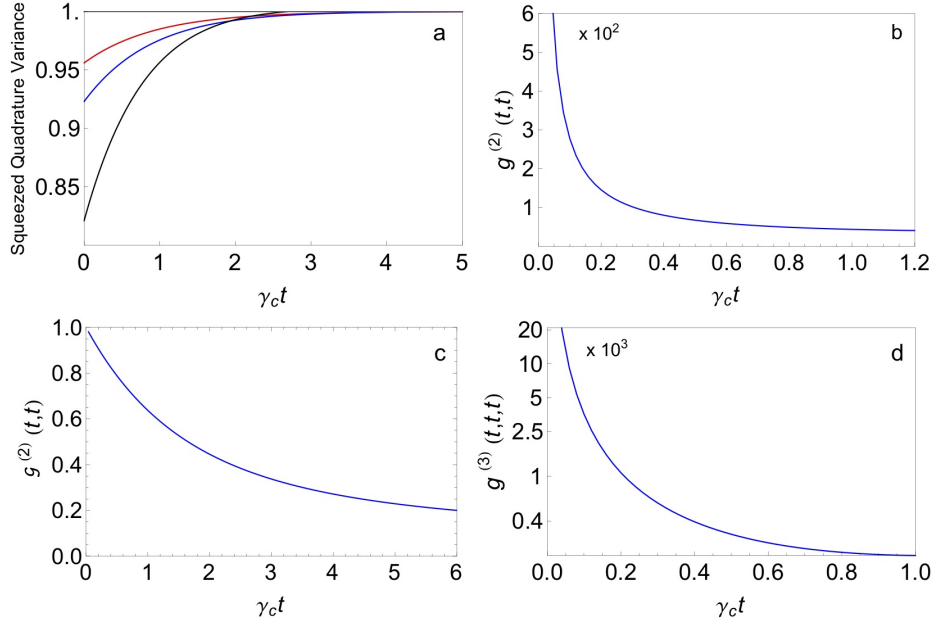


FIG. 3: **Photon statistics.** a: Variance of the squeezed quadrature $X = a + a^\dagger$ for different coupling rates (see text). b: Normalized second order correlation function $g^{(2)}(t, t)$. c: $\mathcal{G}^{(2)}(t, t) = \langle X^-(t) X^+(t) \rangle g^{(2)}(t, t)$. d: Normalized third-order correlation function $g^{(3)}(t, t)$.

Numerical results on the second configuration (sketched in Fig. 1b) are shown in Fig. 2d. The time behaviour of the released photon flux is analogous to that obtained in Fig. 2a for the first configuration. The red line has been obtained for $\omega_s = 2\omega_0$ when only the transition $|s, 0\rangle \rightarrow |\tilde{3}\rangle$ is activated ($\omega_{\tilde{3}} < \omega_s < \omega_{\tilde{4}}$). The black line has been obtained for $\omega_s = 2.5\omega_0$ where also the transition $|s, 0\rangle \rightarrow |\tilde{4}\rangle$ can contribute. We used $\Omega_R/\omega_0 = 0.2$ and the other parameters are the same as those adopted for Fig. 2a.

Conspicuous information on the ongoing physics can be obtained studying the statistics of the emitted photons. A numerical calculation of the variance of the quadrature $X = a + a^\dagger$, for different coupling rates ($\Omega_R/\omega_0 = 0.3, 0.4, 0.6$), is shown in Fig. 3a (other parameters are the same as those adopted for Fig. 2a). Squeezing at initial times coincides with the squeezing exhibited by the false vacuum state $|\tilde{0}\rangle$ and increases with increasing Ω_R/ω_0 . Nevertheless in the present case squeezing is associated to observable photons. The equal-time normalized second order correlation function for cavity photons, $g^{(2)}(t) = \langle X^-(t)X^-(t)X^+(t)X^+(t) \rangle / \langle X^-(t)X^+(t) \rangle^2$, shown in Fig. 3b certifies a highly super-Poissonian statistics, evidencing that cavity photons are released in pairs. Such pair correlation can be further confirmed by calculating $\mathcal{G}^{(2)}(t, t) = \langle X^-(t)X^-(t)X^+(t)X^+(t) \rangle / \langle X^-(t)X^+(t) \rangle$. Such correlation function compares the coincidence rate with the ordinary photodetection rate. When photons are emitted in pairs $\mathcal{G}^{(2)}(t, t) \approx 1$. The decay of these correlation functions is due to cavity losses which tend to destroy correlations. Finally we calculated the equal-time third order normalized correlation function, related to the probability of detecting two cavity photons and a spontaneous emitted photon in the external light modes,

$$g^{(3)}(t) = \frac{\langle \sigma^-(t)X^-(t)X^-(t)X^+(t)X^+(t)\sigma^+(t) \rangle}{\langle \sigma^-(t)\sigma^+(t) \rangle \langle X^-(t)X^+(t) \rangle^2}. \quad (4)$$

The obtained huge value of $g^{(3)}(t, t)$ (see Fig. 3d) confirms the emission mechanism outlined above. Figs. 3b-d have been obtained by using $\Omega_R/\omega_0 = 0.6$ and the other parameters as in Fig. 2a.

In conclusion we proposed a mechanism enabling the release of virtual photon pairs which does not require external forces or time dependent parameters. It resembles the production of particles during the early universe expansion induced by the decay of a false vacuum according to inflationary cosmology. An observation of this mechanism would be possible in state of the art experiments with circuit QED and intersubband polaritons. The spontaneous Casimir effect opens the way to direct investigation of the most interesting feature of ultrastrong coupling, namely the cavity quantum electrodynamic ground state. Such *vacuum spectroscopy* is able to measure the weight of virtual photon pairs in a dressed vacuum or the weight of zero photon states in the excited dressed states. Generalizations of our study to arrays of coupled cavities would form interesting perspectives for future research on virtual photons and pair creation in extended systems [29, 30]. In such systems

the analogy with the scalar field of the inflationary model would be even more stringent.

METHODS

The Liouvillian superoperator in the master equation adopted for the numerical calculations, describing the cavity ($c = 0$) and the material system losses $c = e \rightarrow g$, and $g \rightarrow s$ (configuration a) or $s \rightarrow g$ (b), is defined as $\mathcal{L}_c \rho(t) = \sum_{j,k>j} \Gamma_c^{jk} \mathcal{D}[|j\rangle\langle k|] \rho(t)$, with the dissipator $\mathcal{D}[\mathcal{O}] \rho = \frac{1}{2}(2\mathcal{O}\rho\mathcal{O}^\dagger - \rho\mathcal{O}^\dagger\mathcal{O} - \mathcal{O}^\dagger\mathcal{O}\rho)$. Standard dissipators are recovered in the limit $g \rightarrow 0$. The relaxation coefficients $\Gamma_c^{jk} = 2\pi d_c(\Delta_{kj}) \alpha_c^2(\Delta_{kj}) |C_{jk}^c|^2$ depend on the spectral density of the baths $d_c(\Delta_{kj})$ and the system-bath coupling strength $\alpha_c(\Delta_{kj})$ at the respective transition frequency $\Delta_{kj} = \Omega_k - \Omega_j$ as well as on the transition coefficients $C_{jk} = \langle j|(c + c^\dagger)|k\rangle$ [17, 25]. These relaxation coefficients can be interpreted as the full width at half maximum of each $|k\rangle \rightarrow |j\rangle$ transition. For the sake of simplicity we assume spectral densities $d_c(\Delta_{kj})$ and coupling strengths α_c^2 to be constant. Hence the relaxation coefficients reduce to $\Gamma_c^{jk} = \gamma_c |C_{jk}^c|^2$, where γ_c are the standard damping rates of a weak coupling scenario.

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- [1] Nation, P. D., Johansson, J. R., Blencowe, M. P. & Nori, F. *Colloquium : Stimulating uncertainty: Amplifying the quantum vacuum with superconducting circuits. Rev. Mod. Phys.* **84**, 1–24 (2012).
 - [2] Springel, V., Frenk, C. S. & White, S. D. M. The large-scale structure of the universe. *Nature* **440**, 1137–1144 (2006).
 - [3] Moore, G. T. Quantum theory of the electromagnetic field in a variable-length one-dimensional cavity. *J. Math. Phys* **11**, 2679 (1970).
 - [4] Wilson, C. M. *et al.* Observation of the dynamical casimir effect in a superconducting circuit. *Nature* **479**, 376–379 (2011).
 - [5] Günter, G. *et al.* Sub-cycle switch-on of ultrastrong light-matter interaction. *Nature* **479**, 376–379 (2011).
 - [6] Niemczyk, T. *et al.* Circuit quantum electrodynamics in the ultrastrong-coupling regime. *Nat. Phys.* **6**, 772–776 (2010).
 - [7] Todorov, Y. *et al.* Ultrastrong light-matter coupling regime with polariton dots. *Phys. Rev.*

- Lett.* **105**, 196402 (2010).
- [8] Schwartz, T., Hutchison, J. A., Genet, C. & Ebbesen, T. W. Reversible switching of ultra-strong light-molecule coupling. *Phys. Rev. Lett.* **106**, 196405 (2011).
 - [9] Hoffman, A. J. *et al.* Dispersive photon blockade in a superconducting circuit. *Phys. Rev. Lett.* **107**, 053602 (2011).
 - [10] Scalari, G. *et al.* Ultrastrong coupling of the cyclotron transition of a 2d electron gas to a THz metamaterial. *Science* **335**, 1323–1326 (2012).
 - [11] Mukhanov, V. *Physical Foundations of Cosmology* (Cambridge University Press, 2005).
 - [12] Schwinger, J. On gauge invariance and vacuum polarization. *Phys. Rev.* **82**, 664–679 (1951).
 - [13] Hawking, S. W. Particle creation by black holes. *Commun. Math. Phys.* **43**, 199–220 (1975).
 - [14] Unruh, W. G. Notes on black-hole evaporation. *Phys. Rev. D* **14**, 870–892 (1976).
 - [15] Liberato, S. D., Ciuti, C. & Carusotto, I. Quantum vacuum radiation spectra from a semiconductor microcavity with a time-modulated vacuum rabi frequency. *Phys. Rev. Lett.* **98**, 103602 (2007).
 - [16] Casanova, J., Romero, G., Lizuain, I., García-Ripoll, J. J. & Solano, E. Deep strong coupling regime of the Jaynes-Cummings model. *Phys. Rev. Lett.* **105**, 263603 (2010).
 - [17] Beaudoin, F., Gambetta, J. M. & Blais, A. Dissipation and ultrastrong coupling in circuit QED. *Phys. Rev. A* **84**, 043832 (2011).
 - [18] De Liberato, S., Gerace, D., Carusotto, I. & Ciuti, C. Extracavity quantum vacuum radiation from a single qubit. *Phys. Rev. A* **80**, 053810 (2009).
 - [19] Ridolfo, A., Leib, M., Savasta, S. & Hartmann, M. J. Photon blockade in the ultrastrong coupling regime. *Phys. Rev. Lett.* (to appear) (arXiv:1206.0944v1 [quant-ph]).
 - [20] Ashhab, S. & Nori, F. Qubit-oscillator systems in the ultrastrong-coupling regime and their potential for preparing nonclassical states. *Phys. Rev. A* **81**, 042311 (2010).
 - [21] Ciuti, C. & Carusotto, I. Input-output theory of cavities in the ultrastrong coupling regime: The case of time-independent cavity parameters. *Phys. Rev. A* **74**, 033811 (2006).
 - [22] Paz, J. P. & Zurek, W. H. Quantum limit of decoherence: Environment induced superselection of energy eigenstates. *Phys. Rev. Lett.* **82**, 5181–5185 (1999).
 - [23] Ciuti, C., Bastard, G. & Carusotto, I. Quantum vacuum properties of the intersubband cavity polariton field. *Phys. Rev. B* **72**, 115303 (2005).
 - [24] You, J. Q. & Nori, F. Atomic physics and quantum optics using superconducting circuits.

- Nature* **474**, 589–597 (2011).
- [25] Ridolfo, A., Vilardi, R., Di Stefano, O., Portolan, S. & Savasta, S. All optical switch of vacuum rabi oscillations: The ultrafast quantum eraser. *Phys. Rev. Lett.* **106**, 013601 (2011).
 - [26] Niemczyk, A., *et al.* Selection rules in a strongly coupled qubit-resonator system. *arXiv*: 1107.0810 (2011).
 - [27] Carusotto, I., De Liberato, S., Gerace, D. & Ciuti, C. Back-reaction effects of quantum vacuum in cavity quantum electrodynamics. *Phys. Rev. A* **85**, 023805 (2012).
 - [28] Kardar, M. & Golestanian, R. The “friction” of vacuum, and other fluctuation-induced forces. *Rev. Mod. Phys.* **71**, 1233–1245 (1999).
 - [29] Hartmann, M. J., Brandão, F. G. S. L., & Plenio, M. B. Strongly interacting polaritons in coupled arrays of cavities. *Nat. Phys.* **2**, 849–855 (2006).
 - [30] Leib, M. & Hartmann, M. J. Bose-Hubbard dynamics of polaritons in a chain of circuit quantum electrodynamics cavities. *New J. Phys.* **12**, 093031 (2010).

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